

BOUNDARY VALUE PROBLEMS FOR $\nabla^2 u + k^2 u = 0$

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A BOUNDARY VALUE PROBLEM FOR $\nabla^2 u + k^2 u = 0$

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Introduction and summary.

Part I: The solution of $u_{xx} + u_{yy} + k^2 u = 0$, with u (or $\frac{\partial u}{\partial r}$) on a half-line is obtained by a simpler and more direct method than those previously found in the literature. The method applies equally well to a wedge of arbitrary opening. An idea of Magnus (1) is used to simplify the solution when the boundary values can be expanded in a series of Bessel functions.

Part II: The problem of diffraction by an infinite cone with a spherical tip is treated by a method similar to the used in part I. Both point source and plane wave excitations are considered. When the cone opening θ_0 is $\frac{\pi}{2}$, the problem reduces to reflection from a plane with a hemispherical boss. This simpler problem has been solved by Twersky (7); our results agree with his solution.

Part I.

(a) We first determine the Green's function for the problem.

A two-dimensional unit source is located at the point (ρ', θ') in front of a semi-infinite sheet ($\theta = 0$).

We wish to solve the two-dimensional problem

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$$(1) \quad \frac{1}{\rho} \frac{\partial}{\partial \rho} \left(\rho \frac{\partial u}{\partial \rho} \right) + \frac{1}{\rho^2} \frac{\partial^2 u}{\partial \theta^2} + k^2 u = - \frac{\delta(\rho - \rho') \delta(\theta - \theta')}{\rho'}$$

$$u(\rho, 0+) = 0$$

$$u(\rho, 2\pi-) = 0 \quad .$$

u obeys the Sommerfeld radiation condition at infinity

$$\lim_{r \rightarrow \infty} \sqrt{r} \left(\frac{\partial u}{\partial r} - iku \right) = 0$$

(r is measured from an arbitrary fixed origin).

It is immediately clear that the angular eigenfunctions of the homogeneous equation are just $\sin \frac{n}{2} \theta$. This is a complete orthogonal set of functions over $0 < \theta < 2\pi$, satisfying the boundary conditions.

We write $u(\rho, \theta) = \sum_{n=1}^{\infty} a_n(\rho) \sin \frac{n}{2} \theta$ where the $a_n(\rho)$ are to be determined.

We multiply equation (1) by $\sin \frac{n}{2} \theta$ and integrate from 0 to 2π , obtaining:

$$\begin{aligned} \pi a_n''(\rho) + \frac{\pi}{\rho} a_n'(\rho) + k^2 a_n(\rho) + \frac{1}{\rho} \int_0^{2\pi} \sin \frac{n\theta}{2} \frac{\partial^2 u}{\partial \theta^2} d\theta \\ = \frac{-\sin \frac{n}{2} \theta'}{\rho'} \delta(\rho - \rho') \end{aligned}$$

(A priori it is possible that the coefficients $a_n(\rho)$ are no better than $O(\frac{1}{n^2})$ for large n , therefore we may not be able to differentiate twice termwise with respect to θ).

Integrating twice by parts and using the boundary conditions on u and $\sin \frac{n\theta}{2}$, we readily obtain

$$a_n''(\rho) + \frac{a_n'(\rho)}{\rho} + (k^2 - \frac{n^2}{4}) a_n(\rho) = \frac{-\sin \frac{n}{2} \theta'}{\pi \rho'} \delta(\rho - \rho') .$$

Hence

$$a_n(\rho) = \begin{cases} A H_{\frac{n}{2}}^{(1)}(k\rho) & \rho > \rho' \\ B J_{\frac{n}{2}}(k\rho) & \rho < \rho' \end{cases} .$$

This satisfies the radiation condition at infinity and the finiteness condition at the origin.

We now require

$$a_n(\rho' +) = a_n(\rho' -), \quad \frac{da_n}{d\rho}(\rho' +) - \frac{da_n}{d\rho}(\rho' -) = \frac{-\sin \frac{n}{2} \theta'}{\pi \rho'} .$$

Recalling the Wronskian relation:

$$\rho' [H_{\frac{n}{2}}^{(1)}(k\rho') J_{\frac{n}{2}}(k\rho') - H_{\frac{n}{2}}^{(1)}(k\rho') J_{\frac{n}{2}}'(k\rho')] = \frac{2i}{\pi k} .$$

We have:

$$a_n(\rho) = \frac{i}{2} \sin \frac{n}{2} \theta' \begin{cases} H_{\frac{n}{2}}^{(1)}(k\rho) J_{\frac{n}{2}}(k\rho') & \rho > \rho' \\ H_{\frac{n}{2}}^{(1)}(k\rho') J_{\frac{n}{2}}(k\rho) & \rho < \rho' \end{cases}$$

We use the standard abbreviation:

$$a_n(\rho) = \frac{i}{2} \sin \frac{n}{2} \theta' H_{\frac{n}{2}}^{(1)}(k\rho >) J_{\frac{n}{2}}(k\rho <). \quad \text{The symbol } \rho > \text{ means } \begin{cases} \rho & \text{when } \rho > \rho' \\ \rho' & \text{when } \rho < \rho' \end{cases}.$$

Therefore we have the Green's function:

$$(2) \quad u(\rho, \theta) = G(\rho, \theta; \rho', \theta') = \frac{i}{2} \sum_{n=1}^{\infty} \sin \frac{n}{2} \theta \sin \frac{n}{2} \theta' H_{\frac{n}{2}}^{(1)}(k\rho >) J_{\frac{n}{2}}(k\rho <)$$

(Formulas of this general nature, derived by the use of the idea of sources on a Riemann surface, are to be found in A. Sommerfeld (2)).

Expression (2) may be rewritten

$$\begin{aligned} G(\rho, \theta; \rho', \theta') &= \frac{i}{2} \sum_{n=1}^{\infty} \sin n\theta \sin n\theta' H_n^{(1)}(k\rho >) J_n(k\rho <) \\ &+ \frac{i}{2} \sum_{n=0}^{\infty} \sin \frac{2n+1}{2} \theta \sin \frac{2n+1}{2} \theta' H_{\frac{2n+1}{2}}^{(1)}(k\rho >) J_{\frac{2n+1}{2}}(k\rho <) . \end{aligned}$$

It is easy to recognize the first series. Had we carried out a similar analysis for a completely infinite sheet, the Green's function would turn out to be $i \sum_{n=1}^{\infty} \sin n\theta \sin n\theta' H_n^{(1)}(k\rho >) J_n(k\rho <)$.

We have therefore

$$\frac{i}{2} \sum_{n=1}^{\infty} \sin n\theta \sin n\theta' H_n^{(1)}(k\rho >) J_n(k\rho <) = \frac{i}{8} \left\{ H_0^{(1)}(kR) - H_0^{(1)}(kR^*) \right\}$$

where R is the distance between the source and the observation point,

R^* is the distance between the reflected source about $\theta = 0$ and the observation point. This result could have been obtained from the familiar addition theorem for cylinder functions. We may write

$$(3) \quad G(\rho, \theta; \rho', \theta') = \frac{1}{8} \left\{ H_0^{(1)}(kR) - H_0^{(1)}(kR^*) \right\} \\ + \frac{i}{2} \sum_{n=0}^{\infty} \sin \frac{2n+1}{2} \theta \sin \frac{2n+1}{2} \theta' H_{\frac{2n+1}{2}}^{(1)}(k\rho >) J_{\frac{2n+1}{2}}(k\rho <) .$$

We are now in a position to solve the problem of diffraction from a half-plane. Let the incoming field be e^{ikx} propagating towards increasing x .

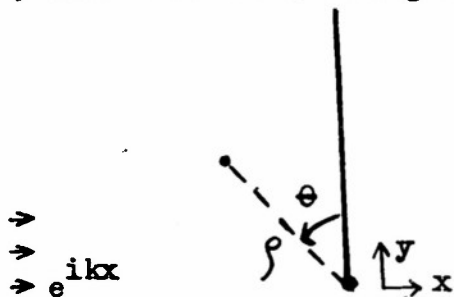


Fig. 1

The half-plane is located in the plane $x = 0$, with $y \geq 0$. We use the coordinate system indicated in Fig. 1. The total field is to vanish for $x = 0$, $y \geq 0$.

We can think of e^{ikx} as being the limiting case of a point source removed to $x = -\infty$.

Indeed suppose there is a two-dimensional unit source at $\theta' = \frac{\pi}{2}$, $\rho = \rho'$.

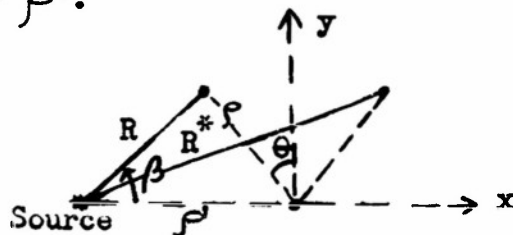


Fig. 2

The field of this point source is just $\frac{i}{4} H_0^{(1)}(kR)$.

We have:

$$R = \frac{(\rho' - \rho \sin \theta)}{\cos \theta} = \frac{\rho' + x}{\cos \theta} .$$

When ρ' is taken very large, the asymptotic expansion for $\frac{i}{4} H_0^{(1)}(kR)$ yields

$$\frac{i}{4} \sqrt{\frac{2}{\pi kR}} e^{-\frac{i\pi}{4}} e^{ikR} .$$

To obtain the limit e^{ikx} as $\rho' \rightarrow \infty$, consider

$$\frac{4}{i} \sqrt{\frac{\pi k \rho'}{2}} e^{\frac{i\pi}{4}} e^{-ik\rho'} \left\{ \frac{i}{4} H_0^{(1)}(kR) \right\} .$$

As $\rho' \rightarrow \infty$, $\cos \beta \rightarrow 1$, $R - \rho' \rightarrow 0$, and it is easy to see that the limit of the above expression is e^{ikx} .

It is reasonable to suppose that performing the same limiting process on (3) will yield the diffraction a plane wave normally incident on a half plane. The result is derived rigorously in (b).

Thus, the total field is given by:

$$u(\rho, \theta) = \lim_{\rho' \rightarrow \infty} \frac{4}{i} \sqrt{\frac{\pi k \rho'}{2}} e^{\frac{i\pi}{4}} e^{-ik\rho'} \left[\frac{i}{8} \left\{ H_0^{(1)}(kR) - H_0^{(1)}(kR^*) \right\} \right. \\ \left. + \frac{i}{2} \sum_{n=0}^{\infty} \sin \frac{2n+1}{2} \theta \sin \frac{2n+1}{2} \frac{\pi}{2} H_{\frac{2n+1}{2}}^{(1)}(k\rho') J_{\frac{2n+1}{2}}(k\rho') \right] .$$

The terms outside the summation yield $\frac{e^{ikx} - e^{-ikx}}{2}$.

In the summation, $\rho < \rho'$ and the asymptotic expansion for $H_{\frac{2n+1}{2}}^{(1)}(k\rho')$ is used:

$$H_{\frac{2n+1}{2}}^{(1)}(k\rho') \sim \sqrt{\frac{2}{\pi k \rho'}} e^{ik\rho'} e^{-in\pi} .$$

We obtain

$$(4) \quad u(\rho, \theta) = i \sin kx + 2e^{-\frac{i\pi}{4}} \sum_{n=0}^{\infty} (i)^{-n} \sin \frac{2n+1}{2} \theta \sin \frac{2n+1}{2} \frac{\pi}{2} J_{\frac{2n+1}{2}}(k\rho) .$$

This formula can easily be identified with that of Magnus (1) who uses the time dependence $e^{i\omega t}$ and a rotated coordinate system.

The results for the case where the normal derivative vanishes on the half-line are easily found:

$$(3a) \quad G(\rho, \theta; \rho', \theta') = \frac{i}{8} \left\{ H_0^{(1)}(kR) + H_0^{(1)}(kR^*) \right\} \\ + \frac{1}{2} \sum_{n=0}^{\infty} \cos \frac{2n+1}{2} \theta \cos \frac{2n+1}{2} \theta' H_{\frac{2n+1}{2}}^{(1)}(k\rho >) J_{\frac{2n+1}{2}}(k\rho <)$$

$$(4a) \quad u(\rho, \theta) = \cos kx + 2e^{-\frac{i\pi}{4}} \sum_{n=0}^{\infty} (i)^{-n} \cos \frac{2n+1}{2} \theta \cos \frac{2n+1}{2} \frac{\pi}{2} J_{\frac{2n+1}{2}}(k\rho) .$$

It is clear that the case of oblique incidence is easily handled by the same method.

(b) Consider the following two-dimensional boundary value problem:

$$\nabla^2 u + k^2 u = 0 \quad u(\rho, \theta+) = u(\rho, 2\pi-) = \text{given function } f(\rho);$$

u is continuous throughout the plane; $\text{grad } u$ is continuous everywhere except possibly on the half-line $\theta = 0$; u obeys a suitable condition at infinity.

Remark on the condition at infinity.

The condition at infinity is not the Sommerfeld radiation condition

$$\lim_{\rho \rightarrow \infty} \sqrt{\rho} \left(\frac{\partial u}{\partial \rho} - iku \right) = 0 .$$

In the case of diffraction by an obstacle having some infinite dimension, neither the total field nor the scattered field can be expected to satisfy the radiation condition.

As a simple example, consider total reflection of a plane wave by an infinite sheet.

The incident field is e^{ikx} .

$$\text{The total field is } \begin{cases} e^{ikx} - e^{-ikx} & x < 0 \\ 0 & x > 0 \end{cases} .$$

The "scattered" field is $-e^{ik|x|}$. None of these fields obeys the Sommerfeld condition. The same is true for diffraction by a semi-infinite sheet.

This is not fatal, however, since all we really need is

$$(5) \quad \lim_{\rho \rightarrow \infty} \int_S \left\{ H_0^{(1)}(kR) \frac{\partial u}{\partial n} - u \frac{\partial}{\partial n} H_0^{(1)}(kR) \right\} d\sigma = 0 .$$

The "scattered" field in the problem of total reflection obeys condition (5).

In general, the geometrical optics field (i.e., no diffraction) will satisfy (5) (see D. S. Jones (3)). Since the scattered field is expected to differ from the geometrical optics field only by a field satisfying the Sommerfeld condition, (5) will hold for the scattered field of the half-plane problem. In any event we adopt (5) as the imposed condition at infinity. It is clear that if (5) holds with the free space Green's function $H_0^{(1)}$, it would also hold with the Green's function G of part I.

Let us apply Green's theorem to u and the Green's function G given by (2). As the boundary of the region of integration is removed to infinity, we obtain (in virtue of (5)):

$$\begin{aligned} u(\rho', \theta') &= \int_0^\infty d\rho \frac{f(\rho)}{\rho} \left\{ \frac{\partial u}{\partial \theta} \Big|_{\theta=0+} - \frac{\partial u}{\partial \theta} \Big|_{\theta=2\pi-} \right\} \\ &= \int_0^\infty \frac{d\rho}{\rho} f(\rho) \left\{ \sum_{n=0}^\infty \frac{1}{2} (2n+1) \sin \frac{2n+1}{2} \theta' H_{\frac{2n+1}{2}}^{(1)}(k\rho >) \right\} . \end{aligned}$$

Under suitable conditions, this may be written as

$$(6) \quad u(\rho', \theta') = \frac{1}{2} \sum_{n=0}^\infty (2n+1) \sin \frac{2n+1}{2} \theta' \int_0^\infty d\rho \frac{f(\rho)}{\rho} H_{\frac{2n+1}{2}}^{(1)}(k\rho >) J_{\frac{2n+1}{2}}(k\rho <) .$$

At first, it might appear that $\lim_{\theta' \rightarrow 0+} u(\rho', \theta') = 0$; this is not true.

We will show that the integral in (6) is no better than $O(\frac{1}{n})$, hence the coefficient of $\sin \frac{2n+1}{2} \theta'$ is no better than $O(\frac{1}{n})$; the series will exhibit discontinuities at $\theta' = 0, 2\pi$, and $\lim_{\theta' \rightarrow \begin{cases} 0+ \\ 2\pi- \end{cases}} u(\rho', \theta') = f(\rho')$.

Formula (6) gives the solution of our problem in a rather untractable form. Magnus (1) noted that integrals of the type (6) are considerably simplified when $f(\rho)$ admits a series development in Bessel functions of integral order. Let

$$I_{m,n} = \int_0^\infty \frac{J_m(k\rho)}{\rho} H_{\frac{2n+1}{2}}^{(1)}(k\rho >) J_{\frac{2n+1}{2}}(k\rho <) d\rho .$$

We split the integral into integrals from 0 to ρ' and ρ' to ∞ .

In each integral we let $t = k\rho$

$$I_{m,n} = \frac{t \left\{ J_m(t) J_{\frac{2n+1}{2}}(t) - J_{\frac{2n+1}{2}}(t) J_m(t) \right\}}{m^2 - \left(\frac{2n+1}{2}\right)^2} \Bigg|_0^{k\rho'} H_{\frac{2n+1}{2}}^{(1)}(k\rho') \\ + \frac{t \left\{ J_m(t) H_{\frac{2n+1}{2}}^{(1)}(t) - H_{\frac{2n+1}{2}}^{(1)}(t) J_m(t) \right\}}{m^2 - \left(\frac{2n+1}{2}\right)^2} \Bigg|_{k\rho'}^{\infty} J_{\frac{2n+1}{2}}(k\rho') .$$

The contribution from 0 vanishes, the contribution at infinity is

$$\frac{2i}{\pi} \frac{e^{\frac{i\pi}{2} \left(m - \frac{2n+1}{2}\right)}}{m^2 - \left(\frac{2n+1}{2}\right)^2} .$$

If we use in addition the formula for the Wronskian of part I, we obtain

$$I_{m,n} = -\frac{2i}{\pi} \frac{e^{-\frac{i\pi}{4}} (i)^{m-n}}{m^2 - \left(\frac{2n+1}{2}\right)^2} J_{\frac{2n+1}{2}}(k\rho') + \frac{2i}{\pi} \frac{J_m(k\rho')}{m^2 - \left(\frac{2n+1}{2}\right)^2} .$$

For the case $f(\rho) = J_m(\rho)$ (6) reduces to:

$$u_m(\rho', \theta') = \frac{(i)^m e^{-\frac{i\pi}{4}}}{\pi} \sum_{n=0}^{\infty} (i)^{-n} \frac{(2n+1) \sin \frac{2n+1}{2} \theta'}{m^2 - \left(\frac{2n+1}{2}\right)^2} J_{\frac{2n+1}{2}}(k\rho') \\ - \frac{J_m(k\rho')}{\pi} \sum_{n=0}^{\infty} \frac{2n+1}{m^2 - \left(\frac{2n+1}{2}\right)^2} \sin \frac{2n+1}{2} \theta' .$$

The first of these Fourier series in θ' has coefficients $O(\frac{1}{n^2})$ due to the rapid decrease with n of $J_{\frac{2n+1}{2}}(k\rho')$. This series represents a continuous function which vanishes at $\theta' = 0, 2\pi$. Hence the limiting values as $\theta' \rightarrow 0+, 2\pi -$ are also zero.

The second series has coefficients $O(\frac{1}{n})$ and may exhibit a discontinuity at $\theta' = 0$. As a matter of fact it can be shown by an elementary computation that

$$-\pi \cos m\theta = \sum_{n=0}^{\infty} \frac{2n+1}{m^2 - (\frac{2n+1}{2})^2} \sin \frac{2n+1}{2} \theta' .$$

Thus:

$$\lim_{\theta' \rightarrow \begin{cases} 0+ \\ 2\pi - \end{cases}} u_m(\rho', \theta') = J_m(k\rho')$$

as required. We can write:

$$u_m(\rho', \theta') = \frac{(i)^m e^{-\frac{i\pi}{4}}}{\pi} \sum_{n=0}^{\infty} (i)^{-n} \frac{(2n+1) \sin \frac{2n+1}{2} \theta'}{m^2 - (\frac{2n+1}{2})^2} J_{\frac{2n+1}{2}}(k\rho') + J_m(k\rho') \cos m\theta'.$$

In general if $f(\rho)$ admits an expansion $f(\rho) = \sum_{m=0}^{\infty} \beta_m J_m(k\rho)$ we will obtain

$$(7) \quad u(\rho', \theta') = \sum_{m=0}^{\infty} \beta_m J_m(k\rho') \cos m\theta' + \frac{e^{-\frac{i\pi}{4}}}{\pi} \sum_{m,n=0}^{\infty} \beta_m (i)^{m-n} (2n+1) \frac{\sin \frac{2n+1}{2} \theta'}{m^2 - (\frac{2n+1}{2})^2} J_{\frac{2n+1}{2}}(k\rho') .$$

It is now easy to solve the problem of diffraction by a half-plane. For simplicity we take normal incidence. The total field $W(\rho, \theta)$ may be split up in an incident field e^{ikx} and a scattered field $u(\rho, \theta)$ satisfying (5). The values of $u(\rho, \theta)$ on the half-plane are just $f(\rho) = -1$. Since $-1 = -\sum_{p=1}^{\infty} \epsilon_p J_{2p}(k\rho)$, we can use (7), with $\beta_m = -\epsilon_m (-i)^m \cos m \frac{\pi}{2}$ ($\epsilon_0 = 1$, $\epsilon_m = 2$, $m > 1$), thus obtaining

$$u(\rho', \theta') = \sum_{m=0}^{\infty} \epsilon_m J_{2m}(k\rho') \cos 2m\theta' + \sum_{n=0}^{\infty} \frac{e^{-\frac{i\pi}{4}}}{\pi} (i)^{-n} (2n+1) \sin \frac{2n+1}{2} \theta' J_{\frac{2n+1}{2}}(k\rho') \sum_{m=0}^{\infty} -\epsilon_m \frac{\cos m \frac{\pi}{2}}{m^2 - (\frac{2n+1}{2})^2}.$$

The first sum is recognized as $-\cos(k\rho' \sin \theta') = -\cos kx'$. The last sum can be evaluated by an elementary computation (see Magnus (1), p. 175)

$$\sum_{m=0}^{\infty} -\epsilon_m \frac{\cos m \frac{\pi}{2}}{m^2 - (\frac{2n+1}{2})^2} = \frac{2}{2n+1} \sin(\frac{2n+1}{2} \frac{\pi}{2}).$$

Thus we have a single summation for $u(\rho', \theta')$ and we obtain

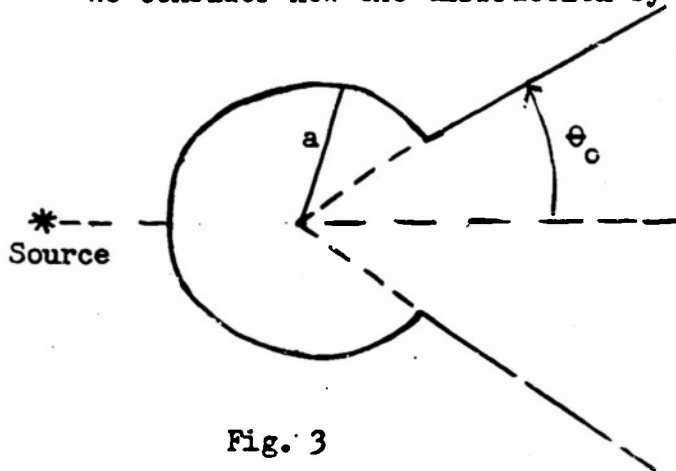
$$(8) \quad W(\rho, \theta) = i \sin kx + 2e^{-i \frac{\pi}{4}} \sum_{n=0}^{\infty} (i)^{-n} \sin \frac{2n+1}{2} \theta \sin \frac{2n+1}{2} \frac{\pi}{2} J_{\frac{2n+1}{2}}(k\rho)$$

which agrees with formula (4) of (a).

It has been shown that (8) can be transformed into the Sommerfeld solution (see Magnus (1), p. 179) and into the Kontorowich-Lebedev solution ((4), p. 241).

Part II.

We consider now the diffraction by a cone with a spherical tip.



The problem of the cone without a tip has been treated by Sollfrey (5).

Fig. 3

The cone is infinite and the source is taken at $\theta = \pi$, $r = r'$. The problem can be solved with only slight additional difficulty for other locations of the source.

The total field satisfies

$$\nabla^2 u + k^2 u = \frac{-\delta(r - r') \delta(\theta)}{2\pi r^2 \sin \theta} \quad (\text{the field is independent of } \varphi)$$

The boundary condition is taken as $u = 0$ on both the sphere and cone.

$$(9) \quad \frac{1}{r^2} \frac{\partial}{\partial r} \left(r^2 \frac{\partial u}{\partial r} \right) + \frac{1}{r^2 \sin \theta} \frac{\partial}{\partial \theta} \left(\sin \theta \frac{\partial u}{\partial \theta} \right) + k^2 u = \frac{-\delta(r - r') \delta(\theta)}{2\pi r^2 \sin \theta}.$$

When $r \neq r'$, the homogeneous equation can readily be separated yielding the equation for the θ part of the separable solutions

$$\frac{1}{\sin \theta} \frac{\partial}{\partial \theta} (\sin \theta w') + \lambda w = 0 \quad \begin{array}{l} w(\theta_0) = 0 \\ w(\pi) \text{ finite} \end{array}$$

λ is a separation constant. If we let $\lambda = \nu(\nu+1)$, $x = \cos \theta$,

$$\frac{d}{dx} [(1-x^2) \frac{dw}{dx}] + \nu(\nu+1)w = 0 \quad \begin{array}{l} w(-1) = \text{finite} \\ w(\cos \theta_0) = 0 \end{array}$$

The only solution which is finite at -1 is $P_\nu(-x)$, the Legendre function of the first kind.

The eigenvalues ν are determined from the equation $P_\nu(-\cos \theta_0) = 0$. The functions $P_\nu(-\cos \theta)$ form a complete orthogonal set in the interval $\theta_0 < \theta < \pi$, with weight $\sin \theta$. The normalization factor

$$\int_{\theta_0}^{\pi} P_\nu^2(-\cos \theta) \sin \theta d\theta = A_\nu$$

will be evaluated later.

We return to equation (9) and consider for u an expansion of the form

$$u = \sum_{\nu} a_{\nu}(r) P_{\nu}(-\cos \theta)$$

We multiply (9) by $r^2 P_{\mu}(-\cos \theta) \sin \theta$ and integrate from θ_0 to π .

We obtain the differential equation for $a_{\mu}(r)$:

$$(10) \quad \frac{d}{dr} (r^2 a'_{\mu}) + k^2 r^2 a_{\mu} - \mu(\mu+1) a_{\mu} = \frac{-\delta(r-r')}{2\pi A_{\nu}}$$

The solution of (10) may be written as:

$$a_{\mu}(r) = C \frac{H_{\mu+1/2}(kr>)}{\sqrt{r>}} \frac{Z_{\mu+1/2}(kr<)}{\sqrt{r<}}$$

where $Z_{\mu+1/2}$ is a cylinder function which vanishes at $r = a$. We have taken care of the boundary conditions at $r = a$, $r = \infty$, and of continuity at $r = r'$. We determined C from the jump condition

$$a'_{\mu}(r'+) - a'_{\mu}(r'-) = - \frac{1}{2\pi r'^2}$$

$$a'_{\mu}(r'+) = C \frac{Z_{\mu+1/2}(kr')}{\sqrt{r'}} \cdot \left\{ \frac{k}{\sqrt{r'}} H_{\mu+1/2}^{(1)}(kr') - H_{\mu+1/2}^{(1)}(kr') \frac{1}{2r'^{3/2}} \right\}$$

$$a'_{\mu}(r'-) = C \frac{H_{\mu+1/2}^{(1)}(kr')}{\sqrt{r'}} \cdot \left\{ \frac{k}{\sqrt{r'}} Z_{\mu+1/2}(kr') - Z_{\mu+1/2}(kr') \frac{1}{2r'^{3/2}} \right\} .$$

Hence

$$C \frac{k}{r'} \text{Wronskian} \left\{ Z_{\mu+1/2}, H_{\mu+1/2}^{(1)} \right\} = - \frac{1}{2\pi r'^2 A_{\nu}} .$$

We write

$$Z_{\mu+1/2}(kr) = H_{\mu+1/2}^{(1)}(kr) - \frac{H_{\mu+1/2}^{(1)}(ka)}{H_{\mu+1/2}^{(2)}(ka)} H_{\mu+1/2}^{(2)}(kr) .$$

It is known that

$$\text{Wronskian} \left\{ H_{\nu}^{(2)}(kr), H_{\nu}^{(1)}(kr) \right\} = - \frac{4i}{\pi kr} .$$

Hence:

$$C = \frac{1}{8A_\nu} \frac{H_{\nu+1/2}^{(2)}(ka)}{H_{\nu+1/2}^{(1)}(ka)}.$$

We may now write the total field:

$$u(r, \theta) = \frac{1}{8} \sum_{\nu} \frac{P_{\nu}(-\cos \theta)}{A_{\nu}} \frac{H_{\nu+1/2}^{(2)}(ka)}{H_{\nu+1/2}^{(1)}(ka)} \frac{H_{\nu+1/2}^{(1)}(kr >)}{\sqrt{r >}} \frac{Z_{\nu+1/2}(kr <)}{\sqrt{r <}}.$$

The sum is taken over all ν such that $P_{\nu}(-\cos \theta_0) = 0$, and $Z_{\nu+1/2}$ is defined by (3).

We now consider the diffraction of a plane wave by our obstacle. This is to be obtained by removing the point source to infinity in a suitable manner. Using the symbols in figure 2 (with $\rho' = r'$), and letting the source go to $-\infty$ along the x axis:

$$\frac{e^{ikR}}{4\pi R} \sim \frac{e^{ikr'}}{4\pi r'} e^{ikx}.$$

The total field $v(r, \theta)$ from a plane wave e^{ikx} will be:

$$\lim_{r' \rightarrow \infty} \frac{4\pi r'}{e^{ikr'}} u(r, \theta)$$

$$= \lim_{r' \rightarrow \infty} \frac{4\pi r'}{e^{ikr'}} \sum_{\nu} \frac{P_{\nu}(-\cos \theta)}{A_{\nu}} \frac{H_{\nu+1/2}^{(2)}(ka)}{H_{\nu+1/2}^{(1)}(ka)} \frac{Z_{\nu+1/2}(kr)}{\sqrt{r}} \frac{H_{\nu+1/2}^{(1)}(kr')}{\sqrt{r'}}.$$

Asymptotically

$$H_{\nu+1/2}^{(1)}(kr') \sim \sqrt{\frac{2}{\pi kr'}} e^{i(kr' - \frac{\nu\pi}{2} - \frac{\pi}{2})}.$$

The total field is:

$$(12) \quad v(r, \theta) = \sum_{\nu} \frac{P_{\nu}(-\cos \theta)}{A_{\nu}} \frac{H_{\nu+1/2}^{(2)}(ka)}{H_{\nu+1/2}^{(1)}(ka)} \frac{Z_{\nu+1/2}(kr)}{\sqrt{r}} \cdot \sqrt{\frac{\pi}{2k}} e^{-\frac{i\nu\pi}{2}}.$$

A_{ν} may be obtained from (see Sollfrey (5)):

$$\begin{aligned} A_{\nu} &= \int_{\theta_0}^{\pi} P_{\nu}^2(-\cos \theta) \sin \theta d\theta \\ &= -\frac{1}{2\nu+1} \sin^2 \theta_0 \left\{ \frac{\partial}{\partial(-\cos \theta)} P_{\nu}(-\cos \theta) \right\}_{\theta=\theta_0} \left\{ \frac{\partial}{\partial \nu} P_{\nu}(-\cos \theta) \right\}_{\theta=\theta_0}. \end{aligned}$$

In the case $\theta_0 = \frac{\pi}{2}$, ν is determined from $P_{\nu}(0) = 0$. Hence ν is an arbitrary odd integer, $2n+1$

$$A_{2n+1} = \int_{\frac{\pi}{2}}^{\pi} P_{2n+1}^2(-\cos \theta) \sin \theta d\theta = \int_0^1 P_{2n+1}^2(x) dx = \frac{1}{4n+3}.$$

We have

$$v(r, \theta) = \sum_{n=0}^{\infty} (4n+3) P_{2n+1}(-\cos \theta) \frac{H_{2n+3/2}^{(2)}(ka)}{H_{2n+3/2}^{(1)}(ka)} \frac{Z_{2n+3/2}(kr)}{\sqrt{r}} \sqrt{\frac{\pi}{2k}} (-1)^n.$$

This was obtained by Twersky (7) using a combination of an image method and superposition.

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